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## WHERE THE NUCLEAR PIONS ARE

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### Abstract

Three recent experiments, which looked at pionic effects in nuclei have concluded that there are *no excess pions*. This puts into serious question the conventional meson-exchange picture of the nucleon-nucleon interaction. Based on arguments of partial restoration of chiral symmetry with density we propose a resolution to this problem.

# 1 Introduction

In an article entitled “Where are the Nuclear Pions?”, Bertsch *et al* [1] discussed three recent experiments which looked for pionic effects in nuclei. One is a  $(\vec{p}, \vec{n})$  quasielastic polarization transfer experiment at LAMPF which, more or less, directly determines the ratio of spin-longitudinal to spin-transverse response functions in the nucleus. At energies below the quasielastic peak conventional models predict this ratio to be significantly larger than unity, while experiment finds a ratio slightly below one. This puts into focus the strength of the tensor interaction in nuclei. In fact, experiment would suggest that at the measured momentum transfer  $V_{\text{tensor}} \sim 0$  while for free nucleon-nucleon interactions it is large. The second is a new deep inelastic muon scattering experiment [2] which no longer sees a significant enhancement in the EMC ratio  $F_2^A/F_2^D$  for  $0.1 \leq x \leq 0.3$ . Since this region is most sensitive to the virtual pion field in the nucleus it was concluded that there are *no excess pions*. Very recently it was realized that the enhancement of the pion field was strongly overestimated in the past because of incorrectly normalized wave functions. Introducing the proper normalization factors there is no contradiction between a conventional model calculation and the measured EMC ratio anymore in the kinematical regime which is dominated by the pion [3]. However, the normalization factors remove only about half of the discrepancy found between a recent Drell-Yan experiment at Fermi Lab [4] and the theoretical predictions. With appropriate choice of kinematics, this experiment directly probes modifications of the sea quarks in the nucleus and is therefore more sensitive to the pion field than the deep inelastic scattering experiment. A pion excess in nuclei would predict a strong  $A$ -dependence of the Drell-Yan ratio and none was observed. Together with the  $(\vec{p}, \vec{n})$  data this calls into serious question the conventional meson-exchange picture of the nuclear interaction.

The authors of ref. [1] suggest that the answer might be found in the modification of gluon properties in the nucleus, suppressing the pion field at distances below 0.5 fm. In the present communication we shall argue that a reasonable explanation lies elsewhere, namely in the partial restoration of chiral invariance with density. Basically, our explanation will

involve the fact that at finite density the hadronic world is “swelled”. More precisely, masses of hadrons made up out of light up and down quarks decrease with density, all at about the same rate [5] thus:

$$\frac{m_N^*}{m_N} \cong \frac{m_\rho^*}{m_\rho} \cong \frac{m_\omega^*}{m_\omega} \dots \cong \frac{f_\pi^*}{f_\pi}, \quad (1.1)$$

where  $f_\pi$  is the pion decay constant, but here its more relevant meaning is that of the order parameter for chiral symmetry breaking. There are approximate signs in the equalities because these were shown [5] to hold only at mean field level, and loop corrections will enter in higher order. The mass of the pion  $m_\pi$  is exempted from this scaling, because it originates from a higher scale than QCD, possibly the electroweak scale. In fact, from the study of pionic atoms [6] we know that due to many-body effects, the pion mass increases slightly, by  $\sim 5$  MeV, in going to the saturation density,  $\rho_0$ , of nuclear matter.

To proceed, we should recall some recent developments in the description of the nucleon-nucleon potential. Some time ago Thomas [7] showed that the data on the sea quark content of the proton can be used to obtain restrictions on the  $t$ -dependence of the  $\pi NN$  vertex function  $\Gamma_{\pi NN}$ . Frankfurt, Mankiewicz and Strikman [8] extended this analysis, finding that the cut-off,  $\Lambda_\pi$ , in a monopole parameterization of  $\Gamma_{\pi NN}$  should be less than 0.5 GeV. Inclusion of more mesons in the nucleon cloud allowed Hwang, Speth and Brown [9] to raise this value to  $\sim 950$  MeV. Clearly such values are too low to correctly describe the NN-scattering data and the binding properties of the deuteron. In various versions of the Bonn potential,  $\Lambda_\pi$  is typically 1.2-1.3 GeV. In an effort to reconcile the deep-inelastic scattering data on the proton with the two-nucleon properties Holinde and Thomas [10] chose  $\Lambda_\pi = 0.8$  GeV but had to introduce an additional pseudoscalar meson (which they call  $\pi'$ ) of mass 1.2 GeV. Assuming a hard form factor of  $\Lambda_{\pi'} = 2$  GeV, the  $\pi' NN$  coupling constant was adjusted to fit the  $NN$  data. More recently it has been recognized [11] that there are at least two objects with pionic quantum numbers, one the elementary pion and the other the correlated  $(\rho\pi)$ -system coupled to the quantum numbers of the pion. The latter may explain the properties of the Holinde-Thomas  $\pi'$  [10]. We find this scenario quite convincing and shall employ the Holinde-Thomas ( $HT$ )

interaction in our calculations.

## 2 Theoretical Development

The enhancement of the pion field is driven by the longitudinal spin-isospin part of the  $NN$  interaction [12]. Employing the  $HT$  interaction [10] it is given by

$$V_{\parallel}(\mathbf{q}, \omega) = [V_{\pi}(q, \omega) + V_{\pi'}(q, \omega)] \boldsymbol{\sigma}_1 \cdot \hat{\mathbf{q}} \boldsymbol{\sigma}_2 \cdot \hat{\mathbf{q}} \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2. \quad (2.1)$$

where

$$V_{\pi}(q, \omega) = \frac{f_{\pi NN}^2}{m_{\pi}^2} \frac{\Gamma_{\pi}^2(q, \omega) q^2}{\omega^2 - (q^2 + m_{\pi}^2)} \quad (2.2)$$

and similarly for the  $\pi'$  meson. In the nuclear medium, this interaction acquires an additional repulsive contribution, usually expressed by the Migdal parameter  $g'$ :

$$\tilde{V}_{\parallel} = V_{\parallel} + \frac{f_{\pi NN}^2}{m_{\pi}^2} g'_{NN} \boldsymbol{\sigma}_1 \cdot \hat{\mathbf{q}} \boldsymbol{\sigma}_2 \cdot \hat{\mathbf{q}} \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2 \quad (2.3)$$

due to short-range correlations induced by the core for the  $NN$  potential. As Baym and Brown have shown [13]  $g'_{NN}$  receives a significant contribution from  $\rho$ -meson exchange, which generates the spin-isospin transverse interaction,  $V_{\perp}$ . This observation will be important to our discussion. The Migdal parameter  $g'$  can be calculated by a momentum-space convolution of the central part of the spin-isospin interaction ( $V_{\text{central}} = 1/3(V_{\parallel} + 2V_{\perp})$ ) with a two-nucleon correlation function  $g$ :

$$\tilde{V}_{\text{central}}(\mathbf{k}, \omega) = \int \frac{d^3k}{(2\pi)^3} g(\mathbf{k} - \mathbf{q}) V_{\text{central}}(\mathbf{q}, \omega). \quad (2.4)$$

To a good approximation  $g(\mathbf{q}) = (2\pi)^3 \delta(\mathbf{q}) - (2\pi^2) \delta(|\mathbf{q}| - q_c)/q^2$ , where  $q_c$  is of the order of the omega-meson mass ( $q_c = 3.94 \text{ fm}^{-1}$  [14]). The resulting  $g'_{NN}$  is  $\omega$  and  $q$  dependent and is displayed in the static limit ( $\omega = 0$ ) as the full line in Fig. 1. For small  $q$  the value is somewhat lower than those extracted from Gamow-Teller systematics [15] ( $g'_{NN} = 0.7 - 0.8$ ). On the other hand it agrees well with G-matrix calculations [16]. The increase in  $g'_{NN}$  at larger  $q$  arises from the  $\rho$ -meson exchange tensor interaction.

Another important ingredient in the description of the virtual pion field is the strong pionic p-wave coupling of the nucleon to  $\Delta(1232)$  isobar. The corresponding spin-isospin correlations are described via transition potentials of the form (2.1) with suitable modifications for the coupling constants and spin-isospin operators. Also in this case, short-range correlations have to be included. Thies [17] pointed out that  $g'_{N\Delta}(0)$  must be close to the classical Lorentz-Lorenz value  $1/3$  in order to explain the absence of multiple scattering in pion-nucleus interaction. Johnson [18] finds  $g'_{\Delta\Delta}(0) = 0.40 \pm 0.13$  which is also consistent with the classical Lorentz-Lorenz value. Therefore we choose  $q_c = 8.66 fm^{-1}$  so that for the  $NN \rightarrow N\Delta$  and  $N\Delta \rightarrow N\Delta$  transition potentials we reproduce  $g'_{N\Delta}(0) = g'_{\Delta\Delta}(0) = 1/3$ . The resulting momentum dependence is also indicated in Fig. 1. The physical origin of the difference in  $g'_{NN}$ ,  $g'_{N\Delta}$  and  $g'_{\Delta\Delta}$  lies in the role of the Pauli principle as was shown by Delorme and Ericson [19] and by Arima *et al* [20] some time ago. Given  $\tilde{V}_{||}^{NN}$  and the corresponding transition potentials all pionic properties, relevant to the experiments under discussion, can be evaluated in linear response theory within the Random-Phase-Approximation. This is the standard scenario employed by many people.

With dropping masses there are several modifications. In medium, the nucleon acquires an effective mass, the mass that enters into the quasiparticle velocity

$$v_{QP} = \frac{p}{m_N^*}, \quad (2.5)$$

where  $p$  is the quasiparticle momentum. By itself this is not unconventional and it is often incorporated in the standard treatment. The crucial point, as was shown by Brown and Rho [5], is that  $m_N^*$  is related to the chiral order parameter  $f_\pi^*$  as

$$\frac{m_N^*}{m_N} \cong \sqrt{\frac{g_A^* f_\pi^*}{g_A f_\pi}} \quad (2.6)$$

once loop corrections are included which bring in the axial vector coupling constant ( $g_A^*$  denotes the in-medium coupling constant). The scaling relation  $m_\rho^*/m_\rho = f_\pi^*/f_\pi$  then directly links the in-medium mass of the  $\rho$  meson to  $m_N^*$  as

$$\frac{m_\rho^*}{m_\rho} = \sqrt{\frac{g_A^* m_N^*}{g_A^* m_N}}. \quad (2.7)$$

A drop of the rho-meson mass with density will increase the range of the spin-transverse interaction,  $\tilde{V}_\perp$ , as well as its coupling constant. We note that [16]

$$\frac{f_{\rho NN}}{m_\rho} = g_{\rho NN} \frac{(1 + \kappa_V^\rho)}{2m_N}, \quad (2.8)$$

where  $\kappa_V^\rho$  is the anomalous  $\rho$ -meson tensor coupling to the nucleon. In accordance with [5]  $g_{\rho NN}$  will not depend on density which follows naturally if one considers the  $\rho$  meson as the gauge particle of the hidden symmetry [21]. With (2.6) we then obtain that

$$\frac{f_{\rho NN}^*}{f_{\rho NN}} = \sqrt{\frac{g_A}{g_A^*}}, \quad (2.9)$$

the  $f_{\rho NN}^*$  being the in-medium coupling. There will also be a change in the  $\pi NN$  coupling constant  $g_{\pi NN}$ . In analogy to (2.8) we have

$$\frac{f_{\pi NN}}{m_\pi} = \frac{g_{\pi NN}}{2m_N}. \quad (2.10)$$

As noted,  $m_\pi$  changes but little with density, increasing by  $\sim 5$  MeV for  $\rho \sim \rho_0$  [6]. In chiral perturbation theory, a change in  $f_{\pi NN}$  enters first in one-loop calculations, four powers higher in the Weinberg counting rules [22] than the basic pion exchange interaction. Consequently, changes in  $f_{\pi NN}$  with density are expected to be small and we neglect them. Thus, the ratio  $f_{\pi NN}/m_\pi$  is taken not to change with density. From the Goldberger-Treiman relation which can be written as

$$\frac{g_{\pi NN}}{m_N} = \frac{g_A}{f_\pi} = \frac{2f_{\pi NN}}{m_\pi} \quad (2.11)$$

this requires  $g_{\pi NN}^*$  to scale as  $m_N^*$  with density. This will turn out to be quite important for the deep inelastic experiments.

The cut-off parameter  $\Lambda_\pi$ , which determines the extent of the pion in  $\pi$ -nucleon interactions and boson exchange models will also be modified in the medium. We can see this in the following way. Consider the  $\pi NN$  vertex and the mass dispersion in the pion channel. Structures, other than elementary pion, will set in with the correlated  $(\rho, \pi)$ -state (see Fig. 2). The mass  $\Lambda_\pi$  involved in this new structure will be determined by the

integral which involves cutting horizontally through the  $\rho$  and  $\pi$  lines, and putting them on shell:

$$\Lambda_\pi \sim m_\rho + m_\pi + T_{\rho\pi} \quad (2.12)$$

where  $T_{\rho\pi}$  is the summed kinetic energy of the  $\rho$  and  $\pi$ . Carrying out the integral in the dispersion relation self-consistently will involve inserting the  $\pi NN$  vertex function appropriately. This has been done many times, most recently by Janssen *et al* [11]. The net result is that

$$\frac{\Lambda_\pi^*}{\Lambda_\pi} \cong \frac{m_\rho^*}{m_\rho}. \quad (2.13)$$

Whereas the kinetic energies  $T_{\rho\pi}$  in eq. (2.12) would seem to increase  $\Lambda_\pi$ , in fact in our picture there are two pions, the  $\pi$  and  $\pi'$ . The latter, when treated as a correlated  $(\rho, \pi)$  system, coupled to the quantum numbers of the  $\pi$ , has a broad mass distribution starting at  $m_\rho + m_\pi$ . The elementary  $\pi$  and the  $\pi'$  mix pushing the  $\pi$  down and the  $\pi'$  up, in energy. It is clear that the main player in the pion vertex function is the mass of the  $\rho$ -meson.

We do not yet have a detailed description of the  $\pi'$ -meson, but much of its mass must come from that of the  $\rho$ -meson. We take  $m_{\pi'}$  and  $\Lambda_{\pi'}$  to scale as in eq. (1.1), although the scaling of the  $\Lambda_{\pi'}$  has little effect. The scaling of  $f_{\pi' NN}$  is assumed to be the same as for  $f_{\rho NN}$ .

Finally we have to scale the mass of the  $\Delta(1232)$  isobar. It is known, e.g. from inclusive electron scattering experiments, that the mass difference between the nucleon and the  $\Delta$  does not significantly change in the nuclear medium. Therefore we keep this difference constant in our calculations:

$$m_\Delta^* - m_N^* = m_\Delta - m_N. \quad (2.14)$$

From the above discussion, once the density dependence of  $m_N^*$  is known, the medium modification of all the other quantities can be derived. For  $m_N^*$  we shall assume a linear density dependence as

$$\frac{m_N^*(\rho)}{m_N} = 1 - 0.3 \frac{\rho}{\rho_0} \quad (2.15)$$

which is adjusted to a value of 0.7 at saturation density as shown in Fig. 3 (for an extensive discussion of the nucleon effective mass in nuclei see [23]). Brown and Rho [24] find, from the measured isovector exchange current in  $^{209}\text{Bi}$ , a value of 0.75 for  $m_N^*/m_N$  averaged over the nucleus, which is consistent with (2.15).

Whereas  $g_A = 1.26$ , the in-medium coupling is roughly

$$g_A^* \simeq 1 \tag{2.16}$$

as is inferred from the quenching of magnetic moments and  $M1$  transitions, as well as the missing Gamow-Teller strength [25, 15]. The renormalization of  $g_A$  has two sources: a screening through virtual  $\Delta$ -hole states [26] and second-order mixing of nucleonic excitations, chiefly through the tensor force [27]. In our model the density dependence of  $g_A^*$  is determined, however, from the assumption that the ratio  $f_{\pi NN}/m_\pi$  remains fixed. Then using eq. (2.6) as well as the Goldberger-Treiman relation (2.11) gives that

$$g_A^* = g_A \left( \frac{m_N^*}{m_N} \right)^{2/3} \tag{2.17}$$

(Fig. 3). At saturation density this yields  $g_A^*/g_A = 0.99$  in close agreement with (2.16). From eq. (2.7) the density dependence of the  $\rho$ -meson mass is also determined (Fig. 3) and we obtain a value of 0.79 for  $m_\rho^*(\rho_0)/m_\rho$ . This is close to QCD sum rule calculations [28, 29].\*

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\*Hatsuda and Lee [29] give 0.82 as their central value. Chanfray and Ericson [30] have shown that exchange current type processes involving the virtual pion field decrease the quark condensate by a factor 10 – 20% over that of Hatsuda and Lee. Birse and McGovern [31] show that with proper calculation of the symmetry breaking matrix element, those result in an enhancement of chiral symmetry breaking in nuclei.



### 3 The Quasifree Polarization-Transfer in the (p,n) Reaction at 495 MeV

The  $(\vec{p}, \vec{n})$  polarization experiments aim at extracting the ratio of the spin-longitudinal to spin-transverse response function

$$\frac{R_{\parallel}(q, \omega)}{R_{\perp}(q, \omega)} = \frac{\sum_{n \neq 0} |\langle 0 | \mathcal{O}_{\parallel} | n \rangle|^2 \delta(\omega - E_n)}{\sum_{n \neq 0} |\langle 0 | \mathcal{O}_{\perp} | n \rangle|^2 \delta(\omega - E_n)} \quad (3.1)$$

where

$$\mathcal{O}_{\parallel} = \sum_{i=1}^A \boldsymbol{\sigma}_i \cdot \mathbf{q} \tau_i^- e^{i\mathbf{q} \cdot \mathbf{r}_i}; \quad \mathcal{O}_{\perp} = \sum_{i=1}^A \boldsymbol{\sigma}_i \times \mathbf{q} \tau_i^- e^{i\mathbf{q} \cdot \mathbf{r}_i} \quad (3.2)$$

from a combination of spin-transfer coefficients. Measurements have been performed in  $^{12}\text{C}$  and  $^{40}\text{Ca}$  at a fixed angle of  $18^\circ$ , corresponding to a peak three-momentum transfer of  $1.72 \text{ fm}^{-1}$ , so as to maximize the difference between  $R_{\parallel}$  and  $R_{\perp}$  expected from the standard treatment of the response functions. The ratio, however, is found to be essentially unity for projectile energy losses below the quasielastic peak [32]. In our picture of dropping masses this can be largely explained. We will sketch below the recent work of Brown and Wambach [33].

When looking at differences in the two-body interaction that cause a deviation of  $R_{\parallel}/R_{\perp}$  from unity, it is clear that these can only come from the tensor part, since

$$\begin{aligned} V_{\parallel} &= V_{\text{central}} + 2V_{\text{tensor}} \\ V_{\perp} &= V_{\text{central}} - V_{\text{tensor}}. \end{aligned} \quad (3.3)$$

For simplicity, let us consider the  $\pi + \rho$  exchange neglecting form factors (the full  $HT$  interaction leads to very similar conclusions). In the static limit

$$V_{\text{tensor}}(q) = \left\{ -\frac{f_{\pi NN}^2}{m_{\pi}^2} \frac{q^2}{(q^2 + m_{\pi}^2)} + \frac{f_{\rho NN}^2}{m_{\rho}^2} \frac{q^2}{(q^2 + m_{\rho}^2)} \right\} S_{12}(\hat{\mathbf{q}}) \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2. \quad (3.4)$$

where  $S_{12}(\hat{\mathbf{q}}) = \boldsymbol{\sigma}_1 \cdot \hat{\mathbf{q}} \boldsymbol{\sigma}_2 \cdot \hat{\mathbf{q}} - 1/3(\boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2)$ . A way of interpreting the experiment is to say that the tensor force is essentially zero at the momentum transfer  $q = 1.72 \text{ fm}^{-1}$ . We adopt as ratio of coupling constants [16]

$$\frac{f_{\rho NN}^2}{m_{\rho}^2} = 2 \frac{f_{\pi NN}^2}{m_{\pi}^2} \quad (3.5)$$

which, within a few percent, is that given by the HT interaction at the relevant momentum transfer. Since  $m_\pi^2 \ll q^2 \ll m_\rho^2$ , the requirement  $V_{\text{tensor}} = 0$  implies that

$$\frac{f_{\rho NN}^2/m_\rho^2}{f_{\pi NN}^2/m_\pi^2} \left( \frac{q^2}{m_\rho^2} \right) \cong 1. \quad (3.6)$$

Using  $m_\pi$  and  $m_\rho$ , however, from the Particle Data book, this ratio is 0.4 instead of unity. From eqs. (2.9) and (3.6) we find that the condition for zero net tensor force at  $q = 1.72 \text{ fm}^{-1}$  becomes<sup>†</sup>

$$\left( \frac{m_\rho}{m_\rho^*} \right)^4 \frac{g_A}{g_A^*} = 2.5, \quad (3.7)$$

or, assuming  $g_A^* = 1$ ,

$$\left( \frac{m_\rho}{m_\rho^*} \right)^4 = 2. \quad (3.8)$$

This gives

$$m_\rho^*/m_\rho = 0.84 \quad (3.9)$$

which is slightly larger than the value of 0.79 at saturation density, found above.

One should note that the  $(p, n)$  reaction at 500 MeV is strongly surface dominated and therefore the nucleus is probed at a lower density. In the calculations of ref. [33] this was taken into account in a semiclassical description, including distortion effects, which amounts to an averaging of the *local* nuclear matter response functions with a weight factor  $F$  as

$$R_{\parallel, \perp}^A(q, \omega) = \int d^3r R_{\parallel, \perp}^{\text{NM}}(q, \omega; \rho(r)) F(r). \quad (3.10)$$

The weight factor can be evaluated using the Eikonal approximation [34]. The average density  $\langle \rho \rangle = \int d^3r \rho(r) F(r) / \int d^3r F(r)$  turns out to be  $0.35\rho_0$  at which  $m_\rho^*/m_\rho = 0.93$ . It is therefore expected that the tensor interaction is reduced but does not completely vanish. As shown in Fig. 4 this is born out of the more complete calculation using the full *HT* interaction as well as all the medium modifications discussed in sect. 2. The results (full lines) for  $^{40}\text{Ca}$  at  $q = 1.72 \text{ fm}^{-1}$  (left panel) and a more recent measurement for  $^{12}\text{C}$  at  $q = 1.2 \text{ fm}^{-1}$  [36] (right panel) give an improvement for energies below the quasielastic

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<sup>†</sup>The loop correction  $g_A$  was not introduced in ref. [33].

peak as compared to a standard RPA treatment (dashed-dotted lines). Our model should be most reliable, however, near the quasielastic peaks of  $\omega \simeq 82 \text{ MeV}$  for  $q = 1.72 \text{ fm}^{-1}$  and  $\omega \simeq 55 \text{ MeV}$  for  $q = 1.2 \text{ fm}^{-1}$  [36], including a  $Q$ -value of  $\sim 20 \text{ MeV}$  for the  $(p, n)$  reaction. For substantially lower  $\omega$  the finite-nucleus collective excitations have to be treated explicitly. For excitation energies above the quasielastic peak, on the other hand, one should bear in mind that contributions from two-step processes become significant [35]. Double scattering contributions to the spin observables are not well understood at present. For  $q = 2.5 \text{ fm}^{-1}$  Brown and Wambach [33] predicted a ratio below unity, which has been experimentally confirmed very recently by Taddeucci *et al.* [36]. However, at this large momentum transfer our results are very sensitive to the exact scaling behavior of the  $\pi'$  which is presently worked out within the microscopic model of ref. [11]. We therefore postpone the discussion of the  $q = 2.5 \text{ fm}^{-1}$  data until these investigations are finished.

The authors of ref. [36] come to the conclusion that the ratio  $R_{\parallel}/R_{\perp}$  remains small because of an unexpected enhancement of the transverse response rather than a non-enhancement of the longitudinal one. There is no theoretical explanation for such an effect. However, the data analysis strongly depends on the treatment of the distortion which enters in terms of an effective number of neutrons. Because of the short range character of the interaction in the transverse channel it is quite reasonable that the distortion is much larger in this channel than in the longitudinal one. Of course, this would also change the ratio.

Our model contains as an essential ingredient the change of the nucleon effective mass  $m_N^*$  with density. It is useful to single out this effect in order to make contact with the recent relativistic calculations of Horowitz and Piekarewicz [37], since the nucleon effective mass enters much in the same way as in the relativistic treatment. A drop of  $m_N^*$  has two effects: (1) the density of particle-hole states is decreased. This shifts the position of the quasielastic peak to higher energies and broadens it. At the same time the longitudinal and transverse correlations are weakened because of an increase of the

particle-hole energies. (2) more importantly  $g'_{NN}$  is increased. Recall that  $g'_{NN}$  receives a significant contribution from  $\rho$  exchange. At zero density and momentum transfer we find the contribution to be  $g'^{(\rho)}_{NN} = 0.28$  while  $\pi$  and  $\pi'$  contribute with 0.08 and 0.21 respectively. By using eq. (2.8) one can relate the increase of the  $\rho$  contribution to  $m_N^*$ :

$$g'^{(\rho)}_{NN}(\rho) \cong g'^{(\rho)}_{NN}(\rho=0) \left( \frac{m_N}{m_N^*} \right)^2. \quad (3.11)$$

Using  $m_N^*/m_N = 0.7$  appropriate for  $\rho_0$  we find  $g'^{(\rho)}_{NN}(\rho_0) = 0.57$ . This would increase  $g'_{NN}$  to a value of 0.86 at nuclear matter density which is very similar to the  $g'_{NN} = 0.9$  of Horowitz and Piekarewicz although they assign only a value of 0.3 from  $\rho$ -exchange. However, in our model the  $\pi'$  scales in the same way as the  $\rho$  and this brings  $g'_{NN}$  up to 1.08 (see Fig. 5). On the other hand this additional short range correlation due to the  $\pi'$  is canceled or even overcompensated by an increase of the  $\pi'$  contribution to the tensor force. Therefore, compared with the results of Horowitz and Piekarewicz we find a somewhat weaker effect of the scaling on the longitudinal-transverse ratio. Preliminary results by Janssen [11] seem to indicate that the  $\pi'NN$  coupling constant coming out of a more microscopic calculation will be much weaker than the HT value. This would considerably improve our results.

So far only the drop of the nucleon mass entered into our argumentation. As can be seen by comparing the solid and dashed lines in Fig. 4 dropping  $m_\rho$  and the other properties as described in section 2 in addition to  $m_N$  does not have a large effect. One concludes that the quasielastic polarization transfer experiments are most sensitive to the change in the nucleon effective mass. In the deep-inelastic scattering experiments the situation is quite different.

## 4 The (Lack of) EMC and Drell-Yan Effects

We shall discuss here the region of  $0.1 < x < 0.3$  where pion enhancement effects were supposed to be and where they weren't [1]. Below  $x = 0.1$  there is shadowing, an interesting phenomenon in its own right. In any description which fits the shadowing, (see,

*e.g.* the 'reggeized' discussion of Brodsky and Lu [38]) and preserves the momentum sum rule, there will be some small overshoot above  $x = 0.1$ . It is argued that this overshoot concerns valence quarks [39], *i.e.* it enters only into the EMC effect. We also will not discuss the dip in the EMC effect in the region of  $x \sim 0.5 - 0.6$ . There are parameterizations of this dip [40]-[42] in terms of rescaling.

In the region  $0.1 < x < 0.3$  the EMC as well as the Drell-Yan ratio are sensitive to the sea quark distributions in the nucleus. We consider a two phase model where the nucleon is made up of a bare quark core and a second component consisting of virtual meson-baryon pairs. Therefore the structure function  $F_2$  reads:

$$F_2^N(x) = Z_N \{ F_2^{core}(x) + \sum_i (\delta F_2^{B_i/N}(x) + \delta F_2^{M_i/N}(x)) \}. \quad (4.1)$$

Here  $Z_N$  denotes a wave function renormalization constant, on which we comment below. The sum in the bracket may run over all meson-baryon decompositions of the nucleon. In deep-inelastic scattering processes the virtual photon couples to the core as well as to the recoil baryon (described by  $\delta F_2^{B_i/N}$ ) and the meson (described by  $\delta F_2^{M_i/N}$ ). The most important example for the latter case is the Sullivan process [43] where the photon couples to the pion cloud (Fig. 6(a)). The corresponding contribution to the structure function of the nucleon is:

$$Z_N \delta F_2^{\pi/N}(x) = Z_N \int_x^1 dy f^{\pi/N}(y) F_2^\pi\left(\frac{x}{y}\right) \quad (4.2)$$

with

$$f^{\pi/N}(y) = \frac{3}{16\pi^2} g_{\pi NN}^2 y \int_{m_N^2 y^2/(1-y)}^\infty dt t \frac{|\Gamma_{\pi NN}(t)|^2}{(t + m_\pi^2)^2} \quad (4.3)$$

being the probability of finding a pion in the nucleon which carries the plus-momentum fraction

$$y = \frac{p_\pi^0 + p_\pi^3}{m_N}. \quad (4.4)$$

The renormalization constant  $Z_N$  (which is missing in the original paper by Sullivan [43]) normalizes the total probability of finding the nucleon in one of the two phases to unity. The importance of this constant in connection with number sum rules has been shown by

Szczurek and Speth [44]. It is given by

$$Z_N = (1 + \sum_i \int_0^1 dy f^{M_i/N}(y))^{-1}, \quad (4.5)$$

where  $f^{M_i/N}$  is the distribution function for the meson  $M_i$ , analogous to eq. (4.3). Taking into account a large set of processes the authors of ref. [44] find  $Z_N \simeq 0.6$  for a cut-off which roughly corresponds to a monopole form factor with  $\Lambda = 800 \text{ MeV}$ . We adopt this value for our calculations.

In the nuclear medium analogous relations hold. We assume that the structure function of the mesons and the baryon cores remain unchanged whereas the meson distribution functions  $f^{M_i}$  and the corresponding ones for the recoil baryon have to be modified. As shown in ref. [3] the modification of the baryon part can be absorbed in the Fermi motion of the nucleons. Furthermore, in the kinematical region we are interested in ( $x \leq 0.3$ ), the only relevant contribution comes from the pion and we can neglect the change of the other meson distribution functions. Thus the structure function  $F_2$  of a nucleon in the nuclear medium becomes:<sup>‡</sup>

$$F_2^{N/A}(x) = \int_x^A dz f^{N/A}(z) F_2^N\left(\frac{x}{z}\right) + \int_x^A dy (Z_A f^{\pi/A}(y) - Z_N f^{\pi/N}(y)) F_2^\pi\left(\frac{x}{y}\right). \quad (4.6)$$

The function  $f^{N/A}(z)$  in the first integral describes the nucleon distribution due to Fermi motion. Since Fermi motion is not very important at small values of  $x$  the main effect comes from the change of the pion distribution function which gives rise to the second integral. Up to this point everything is like in the conventional pion excess model. However, in eq. (4.6) the pion distribution functions  $f^{\pi/N}$  and  $f^{\pi/A}$  are multiplied by the normalization factors  $Z_N$  and  $Z_A$ , respectively. Assuming that we can neglect the change in the distribution functions for mesons other than the pion it follows from eq. (4.5) and

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<sup>‡</sup>For simplicity we discuss only isospin averaged structure functions. In nuclei with neutron excess, like  $^{56}\text{Fe}$ , the pion cloud contains more  $\pi^-$  than  $\pi^+$  mesons, i.e. a small amount of negative charge is transferred from the nucleons to the pion cloud. This effect is properly taken into account in our numerical calculations although it is almost negligible. For details see ref. [3].

the analogous equation for  $Z_A$ :

$$Z_A^{-1} = Z_N^{-1} + \int_0^1 dy (f^{\pi/A}(y) - f^{\pi/N}(y)), \quad (4.7)$$

*i.e.* the normalization factor  $Z_A$  decreases with an increasing distribution function  $f^{\pi/A}$ . Thus the second integral in eq. (4.6) is much smaller (about a factor of  $Z_N^2$ ) than it would be without normalization factors.

In our model the in-medium distribution function is given by

$$f^{\pi/A}(y) = \frac{3}{16\pi^2} g_{\pi NN}^2 y \int_{m_N^2 y^2}^{\infty} dt \int_0^{\frac{t-m_N^2 y^2}{2m_N y}} d\omega t \frac{|\Gamma_{\pi NN}(t)|^2 R_{\parallel}(\omega, \sqrt{t})}{(t + m_{\pi}^2)^2}, \quad (4.8)$$

with  $R_{\parallel}(\omega, q)$  being the non-relativistic spin-longitudinal response function for nuclear matter. It describes Pauli blocking as well as rescattering corrections (see Fig. 6(b)). Pauli blocking leads to a depletion in the mean number of pions per nucleon as compared to the free nucleon which is overcompensated by the rescattering, chiefly through  $\Delta$ -hole excitations, giving a net pion excess.

As in eq. (4.3),  $t$  is the (space-like) four-momentum transfer  $t = \vec{q}^2 - \omega^2$ . The integration limits for  $t$  and  $\omega$  follow directly from eq. (4.4) (with  $p_{\pi}^0 = -\omega$  and  $|\vec{p}_{\pi}| = |\vec{q}|$ ). This is different from the distribution functions which can be found in the literature [46] [47], where the three-momentum transfer is the integration variable (leading to  $\omega_{max} = |\vec{q}| - m_N y$ ) as well as the second argument of  $R_{\parallel}$ . In the non-relativistic regime, *i.e.* when the main contributions to the integral come from regions with  $\omega \ll |\vec{q}|$ , both prescriptions become identical. In addition, however, eq. (4.8) has the correct relativistic low density limit:

$$\lim_{k_F \rightarrow 0} f^{\pi/A}(y) = f^{\pi/N}(y), \quad (4.9)$$

which is not the case for the function given in refs. [46] and [47]. Since both functions,  $f^{\pi/A}$  and  $f^{\pi/N}$ , enter into eq. (4.6) we prefer the more consistent prescription eq. (4.8). In our final results this enhances the pion contribution by a few percent.

For the Fe nucleus we choose an average density of  $\langle \rho \rangle = 0.87\rho_0$  corresponding to

$k_F = 260$  MeV [48]<sup>§</sup>. With no medium modified masses we obtain for the first moment of the pion distribution function  $M_1^{\pi/A} = \int dy f^{\pi/A}(y)$  a value of 0.70 at this density. This has to be compared with the free value of  $M_1^{\pi/N} = 0.41$  (see also Fig. 7). However, because of the normalization factors this enhancement has much less influence on the EMC and Drell-Yan ratios than expected in the past. The corresponding predictions are displayed as the dashed lines in Figs. 8 and 9. In the region  $0.1 \leq x \leq 0.3$  there is no significant deviation of the predicted EMC ratio from the data. Of course, this does not mean that the pion field *is* amplified. Rather, since the more sensitive Drell-Yan data remain strongly overestimated, we are still led to the conclusion that the standard picture fails. This cannot be reconciled by changes in the key parameter  $g'_{N\Delta}$  which would have to be chosen unrealistically large. It also seems implausible that more sophisticated many-body approaches will cure this problem.

With dropping masses, coupling constants and formfactors several modifications occur. Brown, Li and Liu [49] pointed out that the nucleon effective mass  $m_N^*$  rather than  $m_N$  should be used at the soft  $\pi NN$  vertices in Fig. 6. The nucleon mass enters at two places into the derivation of the Sullivan formula (eq. (4.3)). The first place is the spin-isospin current which mixes the large and the small spinor components of the nucleon. Secondly the energy of the virtual pion, as a function of its momentum, is determined from the on-shell condition for the initial and the final nucleon. In both cases the nucleon effective mass should be used.

The modification of the Sullivan formula due to Brown/Rho scaling can be obtained most easily by replacing all properties on the r.h.s. of eq. (4.2) and in eq. (4.3) by the scaled ones. This also includes the variables  $x$  and  $y$  which have to be replaced by  $x^* = \frac{Q^2}{2m_N^*\nu}$  and  $y^* = \frac{p_\pi^0 + p_\pi^3}{m_N^*}$ . Substituting back to the variable  $y = \frac{m_N^*}{m_N} y^*$  we find

$$\delta F_2^{\pi/N^*}(x) = \int_x^{\frac{m_N}{m_N^*}} dy f_\pi^{N^*}(y) F_2^\pi\left(\frac{x}{y}\right) \quad (4.10)$$

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<sup>§</sup>Of course, the Drell-Yan (and EMC) experiments see a higher  $\langle\rho\rangle$  than the polarization transfer, because in the latter case the projectile is affected by the strong interactions.



with

$$f_{\pi}^{N*}(y) = \frac{3}{16\pi^2} g_{\pi NN}^{*2} \left(\frac{m_N}{m_N^*}\right)^2 y \int_{m_N^2 y^2 / (1 - \frac{m_N}{m_N^*} y)}^{\infty} dt t \frac{|\Gamma_{\pi NN}^*(t)|^2}{(t + m_{\pi}^2)^2}. \quad (4.11)$$

Note that the coupling constant  $g_{\pi NN}^*$  comes together with a factor  $\frac{m_N}{m_N^*}$ . As we have argued in sect. 2 this combination is density independent as long as we keep  $f_{\pi NN}/m_{\pi}$  constant. Compared with eq. (4.3), eq. (4.11) leads to a reduced number of pions: The first reason is the smaller cutoff  $\Lambda_{\pi NN}^*$  in the  $\pi NN$  form factor. The second reason is the enhanced lower limit of the  $t$ -integration.

Including nucleon-hole and  $\Delta$ -hole rescattering diagrams (Fig. 6(b)) amplifies the pion field again. Because of the stronger short-range repulsion the effect is smaller than without scaling but it is still present. The general behavior can be seen from Fig. 7 where the first moment of the pion distribution function,  $M_1^{\pi/A}$ , is plotted as a function of density. Whereas  $M_1^{\pi/A}$  is strongly enhanced in the standard RPA calculation (dashed-dotted line) the scaled result (solid line) comes quite close to the free nucleon value (dotted line). At  $\rho = .87\rho_0$  which corresponds to the averaged density of the *Fe* nucleus we find  $M_1^{\pi/A} = 0.70$  without and  $M_1^{\pi/A} = 0.39$  with Brown/Rho scaling which has to be compared with  $M_1^{\pi/N} = 0.41$  for the free nucleon. The dashed line shows the result of the scaled first-order calculation (eq. (4.9)).

As can be expected from these results we almost produce a null effect in the Drell-Yan as well as in the EMC experiments (full lines of Fig. 8 and Fig. 9). This can be seen by comparison with the dotted lines, which show the result of a calculation with the pion contribution (second integral of eq. (4.6)) being switched off. The EMC data in the  $x$ -region of interest are even somewhat underestimated. To obtain the change in the quark distributions and the nucleon structure function (eq. (4.6)) we have employed the quark distributions of the free nucleon and the pion by Owens [45]. Other parameterizations [50]-[52] yield basically the same result. In the EMC calculation nuclear separation energy effects and Fermi motion have been put in as in refs. [49] and [53]. As discussed by Li, Liu and Brown [53], only part of the dip at larger  $x \sim 0.6$  is explained by binding energy effects, once the proper normalization for the baryon number is used. As mentioned

above, the enhancement in the region of  $x = 0.1$  to  $0.2$  seems to be in the valence quarks and such an enhancement can sensibly come from antishadowing or any description of the shadowing which preserves the momentum sum rule.

The pion enhancement in both Drell-Yan and EMC experiments involves convolutions over a fairly wide region of momenta. In this sense, they are less specific than the Los Alamos polarization transfer experiments. While in the former the dropping nucleon mass and the resulting density dependence of  $g'_{NN}$  is the key physical effect, in agreement with the findings by of Horowitz and Piekarewicz [37], the deep inelastic experiments are much more sensitive a change of the  $\rho$ -meson mass, chiefly through the softening of the  $\pi NN$  vertex in the medium. The Drell-Yan data cannot be described if this softening is not taken into account.

## 5 Summary

We have analyzed three recent experiments which have looked at effects of an enhancement of the virtual pion field in nuclei. The negative outcome can be understood from a perspective of partial restoration of chiral symmetry with density which reflects itself in a drop of hadron masses, especially the nucleon and  $\rho$ -meson mass. Following recent developments in the nucleon-nucleon interaction which try to reconcile the sea quark distribution in the nucleon with the low-energy properties of the two-nucleon system, as well as employing microscopic calculations of the  $\pi NN$  form factor we are able to explain the apparent lack of pion enhancement in nuclei. Thus the large discrepancies [1] between the conventional theory and experiment are removed.

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## Figure Captions

**Fig. 1** The momentum dependence of the Fermi liquid parameters  $g'_{NN}$ ,  $g'_{N\Delta}$  and  $g'_{\Delta\Delta}$  at zero density deduced from the potential [10] by using a realistic two-body correlation function. For  $g'_{N\Delta}(0)$  and  $g'_{\Delta\Delta}(0)$  the classical Lorentz-Lorenz value was taken in agreement with the data analysis by Thies [17] and Johnson [18].

**Fig. 2** The vertex function for the  $\pi NN$  interaction. Here we have displayed only the simplest contribution. Because of self interactions, a large number of higher-order diagrams are possible [11].

**Fig. 3** The density dependence of the key quantities in our description. Assuming a linear dependence for  $m_N^*(\rho)/m_N$  such that  $m_N^*(\rho_0)/m_N = 0.7$  the effective  $\rho$ -meson mass is fixed by eq. (2.7) while  $g_A^*(\rho)$  is determined by the requirement that the  $\pi NN$  coupling  $f_{\pi NN}/m_\pi$  in the medium is the same as in free space. The empirical value of  $m_N^*$  is taken from ref. [23] while that of  $g_A^*$  is estimated from [15].

**Fig. 4** The ratio of spin-longitudinal to -transverse response functions in  $^{40}\text{Ca}$  as a function of excitation energy and at fixed momentum transfer  $q = 1.7 \text{ fm}^{-1}$  (left panel). The dashed line gives the result of a standard RPA treatment with only nucleon effective mass, while the full line includes the effects of medium-dependent nucleon mass and meson masses. The dashed-dotted line displays the result without any effective mass. The data (measured at a fixed angle of  $18^\circ$  corresponding to  $q = 1.7 \text{ fm}^{-1}$  at and below the quasielastic peak  $\omega = 80 \text{ MeV}$ ) were taken from refs. [32] (open circles) and [36] (solid circles). The right panel displays the predictions for a momentum transfer of  $^{12}\text{C}$  at  $1.2 \text{ fm}^{-1}$  recently measured at LAMPF [36] ( $\theta = 12.5^\circ$ ). The labeling of the curves is the same as in the left panel.

**Fig. 5** left panel: the density dependence of the Fermi liquid parameters  $g'(0)$  after inclusion of medium-modified masses. The empirical values for  $g'_{NN}$  and  $g'_{\Delta\Delta}$  (see text)

are also given, including their uncertainties.

right panel: the momentum dependence of the Fermi liquid parameters  $g'$  at saturation density,  $\rho_0$ .

**Fig. 6** (a) Deep inelastic scattering off a pion in lowest order. (b) Rescattering correction to (a). The shaded areas in the bubbles are vertex corrections which introduce the local field correction  $g'$ ; nucleon-hole and  $\Delta$ -hole intermediate states are included similarly as in the works of refs. [46] and [47].

**Fig. 7** The first moment  $M_1^{\pi/A} = \int_0^1 dy f^{\pi/A}(y)$  of the pion distribution function as a function of density. The dotted line indicates the free nucleon value  $M_1^{\pi/A} = .41$ . The standard RPA result without Brown/Rho scaling is represented by the dashed-dotted line while the solid line corresponds to the calculation with scaling. The dashed line shows the result of the scaled first-order calculation (no rescattering), corresponding to eq. (4.11).

**Fig. 8** The EMC ratio: The dashed line gives the result of a conventional RPA treatment without medium modifications of the hadron masses, while the full line includes those effects. Wave function renormalization constants have been used in both cases as described in the text. Switching off the pion contribution (second integral of eq. (4.6)) one obtains the result represented by the dotted line. The calculations have been performed at  $\rho = 0.87\rho_0$  which corresponds to the average density of  $^{56}\text{Fe}$ . The data were taken from refs. [2] ( $^{40}\text{Ca}$ ) and [54] ( $^{56}\text{Fe}$ ).

**Fig. 9** The Drell-Yan ratio: The dashed line gives the result of a conventional RPA treatment without medium modifications of the hadron masses, while the full line includes those effects. The dotted line represents the result of a calculation where the pion contribution has been switched off. The data were taken from ref. [4] ( $^{56}\text{Fe}$ ).

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